

# $\pi N$ and $\pi\pi N$ Couplings of the $\Delta(1232)$ and its Chiral Partners

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We investigate the interactions and chiral properties of the four spin- $\frac{3}{2}$ -baryons :  $N^-(D_{13})$ ,  $N^+(P_{13})$ ,  $\Delta^+(P_{33})$  and  $\Delta^-(D_{33})$  together with the nucleon. We construct the  $SU(2)_R \times SU(2)_L$  invariant interactions between the spin- $\frac{1}{2}$  and  $\frac{3}{2}$  baryons with the aid of a new, specially developed spin and isospin projection technique for these baryon fields, where the chiral invariant interactions contain one- and two-pion couplings. We obtain simple relations for the coupling constants of the one- and two-pion spin  $\frac{1}{2} - \frac{3}{2}$  transitions terms. The relation for the one-pion interactions reasonably agrees with the experiments, which suggests that these spin- $\frac{3}{2}$  baryons are chiral partners.

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Chiral symmetry  $SU(2)_R \times SU(2)_L$  is a key property of the strong interaction. When the spontaneous break-down  $SU(2)_R \times SU(2)_L \rightarrow SU(2)_V$  occur, the broken symmetry plays a dynamical role in various scattering processes involving its Nambu-Goldstone bosons, i.e. the pions. Hadrons are then classified by isospin multiplets of the residual symmetry. If chiral symmetry is restored at high temperature or density, hadrons should form degenerate multiplets as classified by the full chiral group representations  $(I_R, I_L)$ , where  $I_R(I_L)$  is isospin for the  $SU(2)_R(SU(2)_L)$  group. Even in the broken world we may expect particles are expressed by one of chiral multiplets or by their simple superpositions [1]. A familiar example is for the chiral mesons  $(\sigma, \vec{\pi})$  and for the vector mesons  $(\vec{\rho}, \vec{a}_1)$ . With regards to the baryons, the role of chiral symmetry in the classification has not been explored much, as one does not know which hadrons would form which chiral multiplets as degenerate partners.

The linear realization of the chiral symmetry offers two merits. First, hadrons in the same chiral multiplet but with different isospins are related by the larger symmetry  $SU(2)_R \times SU(2)_L$  than  $SU(2)_V$ . This can help reducing the number of the free parameters. Second it is easy to investigate the property changes towards the chiral restoration as functions of the chiral condensate. Having these advantages, the purpose of this paper is to investigate the properties of baryons with respecting the chiral symmetry  $SU(2)_R \times SU(2)_L$ , especially we focus on baryon resonances.

It is particularly interesting that the recent studies [2, 3, 4, 5, 6, 7] show that the  $\Delta_{P_{33}}^+(1232)$  and  $N_{D_{13}}^-(1520)$  resonances are qualitatively reproduced in the *quenched* lattice QCD, which validates to some extent the empirical assumption that the baryons are dominated by their  $3q$  Fock components. Recently, we clarified the relation between the chiral multiplets and the quark structures [8]. For instance, interpolating fields

used in Ref. [2, 3, 4, 5, 6, 7] belong to the chiral multiplet  $(1, \frac{1}{2}) \oplus (\frac{1}{2}, 1)$  [8]. This is our starting assumption where a set of spin- $\frac{3}{2}$  baryons form the chiral multiplet as chiral partners. We extend this idea to include two other four-star resonances, the  $N_{P_{13}}^+(1720)$  and  $\Delta_{D_{33}}^-(1700)$ , following Jido et. al. [9], where the four spin- $\frac{3}{2}$  baryons form a certain set of chiral multiplets, so-called Quartet scheme. Ref. [9] mostly discussed the interactions between these chiral multiplets with the same spins, but not with the different spins especially the nucleon. Inclusion of the nucleon enables us to testify such a framework in comparison with the experimental data for not only masses but also other quantities related to the dynamical processes such as resonance decays and scatterings.

In this Letter, we construct an effective Lagrangian for four types of four-star resonances,  $\Delta(1232)$ ,  $N(1520)$ ,  $N(1720)$  and  $\Delta(1700)$  together with the nucleon and investigate the structures of the one- and two-pion coupling strengths. We derive a relation among the one-pion coupling constants of the four baryon resonances, which agree well with the experimental data. We also consider the property changes of the one-pion couplings towards the chiral restoration.

We begin with the nucleon's chiral multiplet for which there are two possibilities  $(\frac{1}{2}, 0) \oplus (0, \frac{1}{2})$  and  $(1, \frac{1}{2}) \oplus (\frac{1}{2}, 1)$  when the nucleon is a three-quark field. We assume the nucleon to be the fundamental representation. There are also two possible chiral representations for the  $\Delta(1232)$ :  $(1, \frac{1}{2}) \oplus (\frac{1}{2}, 1)$  and  $(\frac{3}{2}, 0) \oplus (0, \frac{3}{2})$ . We choose the former case as is commonly used to describe the  $\Delta(1232)$ , as well as in Ref. [9]. Now we define two types of the diquarks; a Lorenz vector iso-scalar diquark  $V^\mu$  ( $I(J)^P = 0(1)^-$ ) and axial-vector iso-vector diquark  $A^{\mu i}$  ( $1(1)^+$ ),

$$V^\mu = \tilde{q}\gamma^\mu q, \quad (1a)$$

$$A^{\mu i} = \tilde{q}\gamma^\mu\gamma_5\tau^i q, \quad (1b)$$

where  $\tilde{q} = q^T C(i\tau_2)\gamma_5$  is a transposed quark field. These diquarks form the chiral multiplet  $(\frac{1}{2}, \frac{1}{2})$ , just like the  $\sigma$  and  $\vec{\pi}$  mesons, which is a key ingredient in our construction of the chiral invariant interactions. We will come

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back to this point later. There is one possible operator for  $I(J) = \frac{3}{2}(\frac{3}{2})$ ,

$$\Delta_4^{\mu i} = A_\nu^j \Gamma_{3/2}^{\mu\nu} P_{3/2}^{ij} q, \quad (2a)$$

and two for  $I(J) = \frac{1}{2}(\frac{3}{2})$ ,

$$N_V^\mu = V_\nu \Gamma_{3/2}^{\mu\nu} \gamma_5 q, \quad (2b)$$

$$N_A^\mu = A_\nu^i \Gamma_{3/2}^{\mu\nu} \tau^i q. \quad (2c)$$

Here the isospin- $\frac{3}{2}$   $P_{3/2}^{ij}$  and the isospin- $\frac{1}{2}$  projection operator  $P_{1/2}^{ij}$ , satisfy  $\delta^{ij} = P_{3/2}^{ij} + P_{1/2}^{ij}$  [8]. Similarly the spin- $\frac{3}{2}$  projection operator  $\Gamma_{3/2}^{\mu\nu}$ , with the spin  $\frac{1}{2}$  projection operator  $\Gamma_{1/2}^{\mu\nu}$ , satisfy the completeness relation  $g^{\mu\nu} = \Gamma_{3/2}^{\mu\nu} + \Gamma_{1/2}^{\mu\nu}$ . Note that with the spin- $\frac{3}{2}$  projection operator the baryon fields still contain four fictitious spin- $\frac{1}{2}$  components. However, the chiral transformation properties does not depend on the choice of the spin projection operators, the local or non-local type [8]. Our strategy is firstly use the local projection operators in the construction of the Lagrangian, later eliminate the spin- $\frac{1}{2}$  components in the calculations of the physical quantities, the one-pion decays in the present context.

Taking into account the normalization and Pauli-principle, which is implemented by the Fierz transformation, we define the baryon fields as

$$\Delta_1^{\mu i} = \frac{\Delta_4^{\mu i}}{2}, \quad (3a)$$

$$N_1^\mu = \frac{\sqrt{3}}{2} \frac{N_V^\mu}{2} + \frac{1}{2} \frac{N_A^\mu}{2\sqrt{3}}, \quad (3b)$$

where we separate the coefficients to show explicitly the normalized baryon fields  $\Delta_4^{\mu i}/2$ ,  $N_V^\mu/2$  and  $N_A^\mu/2\sqrt{3}$ . Note that the mixing between  $N_V^\mu$  and  $N_A^\mu$  results from the chiral transformations of  $V^\mu$  and  $A^{\mu i}$ , and the mixing angle between  $N_V^\mu$  and  $N_A^\mu$  are determined by the Fierz transformation [8]. The chiral transformation properties are straightforwardly given by

$$\delta_5^{\vec{a}} N_1^\mu = \frac{5}{3} i \vec{a} \cdot \vec{\tau} \gamma_5 N_1^\mu + \frac{4}{\sqrt{3}} i \gamma_5 \vec{a} \cdot \vec{\Delta}_1^\mu, \quad (4a)$$

$$\begin{aligned} \delta_5^{\vec{a}} \Delta_1^{\mu i} &= \frac{4}{\sqrt{3}} i \gamma_5 a^j P_{3/2}^{ij} N_1^\mu - \frac{2}{3} i \tau^i \gamma_5 \vec{a} \cdot \vec{\Delta}_1^\mu \\ &+ i \vec{a} \cdot \vec{\tau} \gamma_5 \Delta_1^{\mu i}, \end{aligned} \quad (4b)$$

implying that a set of  $N_1^\mu$  and  $\Delta_1^{\mu i}$  form the chiral multiplet  $(1, \frac{1}{2}) \oplus (\frac{1}{2}, 1)$ .

Even after establishing these chiral transformations, it is a non-trivial task to build chirally invariant interactions from these fields, so we shall develop a new method to project out the good spin and isospin parts from chiral invariant operators containing reducible products of three-quark fields. This projection technique is performed in two steps. Firstly we adopt a multiple product

of meson and baryon operators described as direct products of quark bi-linears (diquarks) and a quark. Here the equivalence of the chiral properties of  $(V^\mu, A^{\mu i})$  and  $(\sigma, \pi^i)$  has been used. Since such composite operators are reducible under both the chiral and the spin and isospin transformations, we perform the decomposition into irreducible parts containing only spin and isospin-projected baryons.

As an illustration let us consider the vector and axial-vector diquarks  $(V^\mu, A^{\mu i})$ . As explained, they belong to the chiral multiplet  $(\frac{1}{2}, \frac{1}{2})$  similar to  $(\sigma, \vec{\pi})$ . Therefore the  $V_\mu^2 + A_\mu^2$  combination is a chiral scalar, which immediately leads to the chiral invariant term  $\bar{q}(V_\mu^2 + A_\mu^2)U_5 q$ , where  $U_5 = \sigma + i\gamma_5 \vec{\tau} \cdot \vec{\pi}$ . The direct products of a quark and diquarks  $V^\mu q$  and  $A^{\mu i} q$  contain several kinds of baryons with  $I(J) = \frac{1}{2}(\frac{1}{2}), \frac{1}{2}(\frac{3}{2}), \frac{3}{2}(\frac{3}{2})$  [14]. The decomposition into irreducible spin and isospin parts is carried out by using the completeness relations of both the spin and isospin projection operators. The resulting interaction Lagrangian is given by

$$\begin{aligned} \mathcal{L}_{\pi BB}^1 &= g_1 \left( \bar{\Delta}_1^i U_5 \Delta_1^{\mu i} - \frac{3}{4} \bar{N}_1 U_5 N_1^\mu \right. \\ &\quad \left. + \frac{1}{12} \bar{N}_1 \tau^i U_5 \tau^i N_1^\mu + \frac{\sqrt{3}}{6} \bar{N}_1 \tau^i U_5 \Delta_1^{\mu i} \right). \end{aligned} \quad (5)$$

Note that the relative weights for  $N_1^\mu$  and  $\Delta_1^{\mu i}$  are unambiguously fixed by Eq. (4) without dependence on any free parameters [15].

Now, following Ref. [9], we introduce a new set of spin- $\frac{3}{2}$  baryons  $(N_2^\mu, \Delta_2^{\mu i})$  that have the  $SU(2)_A$  transformation properties opposite in sign to those of  $(N_1^\mu, \Delta_1^{\mu i})$ , so called mirror baryons [9, 10, 11, 12]. They may appear due to some more complicated structures such as non-local nature of three-quark states and multiquark components. Here we assume the existence of the mirror baryons without considering their internal structures in detail. The diagonal interactions for the mirror baryons are easily obtained:

$$\begin{aligned} \mathcal{L}_{\pi BB}^2 &= g_2 \left( \bar{\Delta}_2^i U_5^\dagger \Delta_2^{\mu i} - \frac{3}{4} \bar{N}_2 U_5^\dagger N_2^\mu \right. \\ &\quad \left. + \frac{1}{12} \bar{N}_2 \tau^i U_5^\dagger \tau^i N_2^\mu + \frac{\sqrt{3}}{6} \bar{N}_2 \tau^i U_5^\dagger \Delta_2^{\mu i} \right). \end{aligned} \quad (6)$$

In addition, the following mass terms are allowed

$$\mathcal{L}_{BB} = -m_0 \left( \bar{\Delta}_1^i \Delta_2^{\mu i} + \bar{N}_1 N_2^\mu \right). \quad (7)$$

In contrast to Eqs. (5) and (6), the mixings between  $N_1^\mu(\Delta_1^{\mu i})$  and  $N_2^\mu(\Delta_2^{\mu i})$  occur only after the mass diagonalization when the so-called mirror mass  $m_0$  is finite [9, 11, 12]. Combining Eqs. (5)~(7), the quartet scheme of Jido et. al. [9] is exactly reproduced.

Next, we include the nucleon and its couplings with the spin- $\frac{3}{2}$  baryons, which is new in this work. Similarly to the above discussion,  $(V^\mu, \vec{A}^\mu)$  and  $(\sigma, \vec{\pi})$  form a

chiral scalar  $\sigma V_\mu + i\boldsymbol{\pi} \cdot \mathbf{A}_\mu$ . Hence we find two chirally invariant interactions: (1)  $\bar{N}U_5[(\partial^\mu \sigma)V_\mu + i(\partial^\mu \boldsymbol{\pi}) \cdot \mathbf{A}_\mu]q$ , and (2)  $\bar{N}(\partial^\mu U_5)(\sigma V_\mu + i\boldsymbol{\pi} \cdot \mathbf{A}_\mu)q$ . With the irreducible decomposition, we obtain

$$\mathcal{L}_{\pi NB}^1 = \frac{g_3}{\Lambda^2} \left[ \bar{N}U_5(i\partial_\mu \pi^i)\Delta_1^{\mu i} + \frac{\sqrt{3}}{2}\bar{N}U_5(\gamma_5\partial_\mu \sigma + \frac{i}{3}\partial_\mu \boldsymbol{\pi} \cdot \boldsymbol{\tau})N_1^\mu \right], \quad (8)$$

$$\mathcal{L}_{\pi NB}^2 = \frac{g_4}{\Lambda^2} \left[ \bar{N}(\partial_\mu U_5)(i\pi^i)\Delta_1^{\mu i} + \frac{\sqrt{3}}{2}\bar{N}(\partial_\mu U_5)(\gamma_5\sigma + \frac{i}{3}\boldsymbol{\pi} \cdot \boldsymbol{\tau})N_1^\mu \right], \quad (9)$$

where the dimensional parameter  $\Lambda$  is introduced to keep the coupling constants dimensionless. We neglect the higher-order terms for the nucleon  $\bar{N}U_5\partial U_5^\dagger N$  and  $\bar{N}(\partial^\mu U_5)U_5^\dagger \gamma_\mu N$ . For the mirror baryons, we obtain the single-meson coupling

$$\mathcal{L}_{\pi NB}^3 = \frac{g_5}{\Lambda} \left( \bar{N}(i\partial^\mu \pi^i)\Delta_{2\mu}^i - \frac{\sqrt{3}}{2}\bar{N}\partial^\mu(\gamma_5\sigma - \frac{1}{3}i\boldsymbol{\tau} \cdot \boldsymbol{\pi})N_{2\mu} \right). \quad (10)$$

Again, we neglect the nucleon term  $\bar{N}\partial U_5^\dagger N_m$ , where  $N_m$  is another nucleon field having the mirror properties. Note that the interactions Eqs. (8) and (9) involve two mesons, while Eq. (10) contains only the single meson couplings.

TABLE I: Masses (second column) and Coupling constants (third column). For masses, we follow Jido et. al [9]. The experimental values are taken from the PDG tables [13]. The experiments determine only the absolute values of the coupling constants, the positive values are our assumption.

States	Masses [MeV]	$g_{\pi NB}/\Lambda$ [MeV <sup>-1</sup> ]	$\Gamma_{B \rightarrow \pi N}$ [MeV]
	Theo (Exp)	Theo (Exp)	
$\Delta_+^{\mu i}$ ( $P_{33}$ )	1320 (1232)	15 (16)	118
$\Delta_-^{\mu i}$ ( $D_{33}$ )	1770 (1700)	9.2 (9.5)	45
$N_-^\mu$ ( $D_{13}$ )	1430 (1520)	9.4 (8.6)	69
$N_+^\mu$ ( $P_{13}$ )	1660 (1720)	2.4 (2.4)	30
<hr/>			
	$m_0 = 1550$	$g_1 = g_2 = 2.4$	
$g_5/\Lambda = 17$	$g_3 f_\pi/\Lambda^2 = 4.2$	$g_4 f_\pi/\Lambda^2 = 8.2$	

Having constructed the Lagrangian with the nucleon and spin- $\frac{3}{2}$  baryons, let us determine the parameters  $g_{1,2}$  and  $m_0$ , following Ref. [9]. The results are shown in Table I [16]. After the diagonalization of the mass term and the corresponding parity (re)definition, the mass eigenstates are obtained as: for the  $\Delta$ s,  $\Delta_+^{\mu i} = (\Delta_1^{\mu i} + \Delta_2^{\mu i})/\sqrt{2}$ ,  $\Delta_-^{\mu i} = \gamma_5(-\Delta_1^{\mu i} + \Delta_2^{\mu i})/\sqrt{2}$ , and for the  $N^*$ s,  $N_-^\mu = \gamma_5(-N_1^\mu + N_2^\mu)/\sqrt{2}$ ,  $N_+^\mu = (N_1^\mu + N_2^\mu)/\sqrt{2}$ , where the subscripts  $\pm$  denote the parity [17].

After the spontaneous breaking  $SU(2)_R \times SU(2)_L \rightarrow SU(2)_V$ , the one-pion interactions in Eqs. (8)-(10) are reduced to

$$\begin{aligned} \mathcal{L}_{\pi NB} = & \frac{g_{\pi N \Delta^+}}{\Lambda} \bar{N}(i\partial_\mu \pi^i)\Delta_+^{\mu i} + \frac{g_{\pi N \Delta^-}}{\Lambda} \bar{N}(i\gamma_5\partial_\mu \pi^i)\Delta_-^{\mu i} \\ & + \frac{g_{\pi N N^{*-}}}{\Lambda} \bar{N}(i\gamma_5\partial_\mu \boldsymbol{\pi} \cdot \boldsymbol{\tau})N_-^\mu \\ & + \frac{g_{\pi N N^{*+}}}{\Lambda} \bar{N}(i\partial_\mu \boldsymbol{\pi} \cdot \boldsymbol{\tau})N_+^\mu, \end{aligned} \quad (11a)$$

where the coupling constants are given by

$$g_{\pi N \Delta^\pm} = \frac{1}{\sqrt{2}\Lambda}(g_5\Lambda \pm g_3f_\pi), \quad (11b)$$

$$g_{\pi N N^{*\pm}} = \frac{\sqrt{6}}{12\Lambda}((g_3 + 3g_4)f_\pi \mp g_5\Lambda). \quad (11c)$$

The three coupling constants  $g_{3,4,5}$  are determined from the one-pion decay widths of the resonances as shown in Table I. We obtain quantitatively reasonable results for all the four coupling constants Eqs. (11). Eliminating  $g_{3,4,5}$  from Eqs. (11), we obtain a new relation:

$$(g_{\pi N \Delta^+} + g_{\pi N \Delta^-}) = 2\sqrt{3}(g_{\pi N N^{*-}} - g_{\pi N N^{*+}}), \quad (12)$$

which satisfies the experimental data with a numerical error of less than 10%. Considering the simplicity of the present description, this is an encouraging result suggesting that the spin- $\frac{3}{2}$  baryons are the candidates of the chiral partners.

One of interesting properties of the present model is the two-pion contact terms, which are an inevitable consequence of the chiral invariance. They involve only the  $g_3$  and  $g_4$ , while  $g_5$ , which is a leading contribution in the one-pion couplings, does not contribute to the two-pion couplings. The two-pion decay of  $\Delta(1232)$  is therefore suppressed by the smallness of the coupling constants as compared to the one-pion decay. On top of this, the derivative coupling causes an additional suppression of the two-pion decay rate, due to the small final state pion momentum. Hence we can expect strong suppression of the two-pion decay of  $\Delta(1232)$ . Explicitly the two-pion contact interactions are given by

$$\begin{aligned} \mathcal{L}_{2\pi NB} = & \frac{1}{\sqrt{2}\Lambda^2} \bar{N}A_\mu^i \Delta_+^{\mu i} - \frac{1}{\sqrt{2}\Lambda^2} \bar{N}A_\mu^i \gamma_5 \Delta_-^{\mu i} \\ & + \frac{\sqrt{6}}{12\Lambda^2} \bar{N}B_\mu \gamma_5 N_-^\mu - \frac{\sqrt{6}}{12\Lambda^2} \bar{N}B_\mu N_+^\mu, \end{aligned} \quad (13a)$$

with

$$\begin{aligned} A_\mu^i = & g_3(i\gamma_5 \boldsymbol{\pi} \cdot \boldsymbol{\tau})(i\partial_\mu \pi^i) + g_4(i\gamma_5 \partial_\mu \boldsymbol{\pi} \cdot \boldsymbol{\tau})(i\pi^i), \quad (13b) \\ B_\mu = & g_3(i\gamma_5 \boldsymbol{\pi} \cdot \boldsymbol{\tau})(i\partial_\mu \boldsymbol{\pi} \cdot \boldsymbol{\tau}) + g_4(i\gamma_5 \partial_\mu \boldsymbol{\pi} \cdot \boldsymbol{\tau})(i\boldsymbol{\pi} \cdot \boldsymbol{\tau}). \end{aligned} \quad (13c)$$

Hence we obtain a relation between the 2- $\pi$  contact terms:

$$|g_{2\pi N \Delta^+}| = |g_{2\pi N \Delta^-}| = 2\sqrt{3}|g_{2\pi N N^{*+}}| = 2\sqrt{3}|g_{2\pi N N^{*-}}|. \quad (14)$$

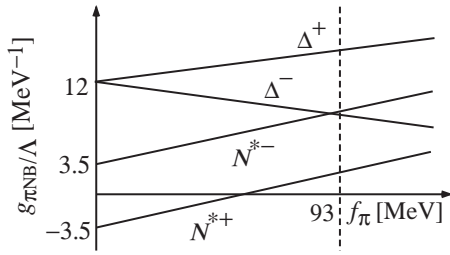


FIG. 1: The illustration of  $f_{\pi}$  dependence of the one-pion coupling constants  $g_{\pi NB}/\Lambda$  in the case (1).

In contrast to the  $\Delta(1232)$  case, it is expected that the two-pion contact term leads to larger contributions for other baryon resonances, because of the larger final state pion momenta. In particular, the two-pion coupling constants of  $N^{*+}(\Delta^-)$  has the same magnitude as compared with that of  $N^{*-}(\Delta^+)$ , while the one-pion coupling constant is suppressed by the negative sign in Eqs. (11). This qualitatively explains the observed feature of the two-pion decay enhancement in the decays of  $N(1720)$  and  $\Delta(1700)$ . Due to the lack of other resonances, such as  $\rho$ -meson and  $N(1440)$ , we do not consider this point in details.

As an application of our results to one of the recent interests, we consider a situation towards the chiral restoration at high temperature or density. We briefly consider the property changes of the one-pion coupling constants in two cases: (1) the scale parameter  $\Lambda$  is constant and (2)  $\Lambda = f_{\pi}$ . In the case (1), as  $f_{\pi}$  decreases, the three coupling constants  $g_{\pi N\Delta^+}$  and  $g_{\pi NN^{*\pm}}$  decrease, while the remaining one  $g_{\pi N\Delta^-}$  increases. At  $f_{\pi} = 0$ , we obtain the simple relation  $g_{\pi N\Delta^+} = g_{\pi N\Delta^-} = \frac{\sqrt{3}}{6} g_{\pi NN^{*-}} = -\frac{\sqrt{3}}{6} g_{\pi NN^{*+}}$ , which is shown in Fig. 1. In the case (2),

all the coupling constants simply increase proportional to  $f_{\pi}^{-1}$ . Eq. (12) does not depend on the value of  $\Lambda$ , hence it always holds in both cases.

In summary, we have investigated the chiral properties of four spin- $\frac{3}{2}$  baryon resonances together with the nucleon. We have constructed the effective interactions for the spin  $\frac{1}{2} - \frac{3}{2}$  transition terms with the aid of the spin and isospin projection formalism for the baryon fields comprised of three quark fields. Of course, we can prove the chiral invariance of the derived interactions directly from the chiral transformation laws, but the results are general from the group-theoretical point of view. Within the  $J = \frac{3}{2}$  sector, the projection formalism reproduces the Quartet scheme proposed by Jido et al. [9]. In addition, we derived the minimal chiral invariant one- and two- meson couplings with spin  $\frac{1}{2} - \frac{3}{2}$  baryons. We found that the one-pion couplings describing the spin  $\frac{1}{2} - \frac{3}{2}$  transitions are constrained by the chiral symmetry via the Eq. (12), which quantitatively agrees with the experiment. Considering the simplicity of our assumptions on the effective Lagrangian, it is an remarkable result suggesting that these baryons are chiral partners. In addition, we obtain chiral two-pion couplings, whose strengths are entirely determined by the one-pion coupling constants. This enable us to predict two-pion decays of the resonances that can be tested in experiments. In this Letter we employed a new projection technique to derive the effective chiral interaction Lagrangians between baryons of different spin and isospin.

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  - [14]  $I(J) = \frac{3}{2}(\frac{1}{2})$  state is forbidden by the Pauli-principle [8] when it is described as a local three quark state.
  - [15] Eq. (5) predicts a mass relation  $m(\Delta_1^{\mu i}) : m(N_1^{\mu}) = 2 : 1$ . There are no candidates for this sort of  $N^*$  and  $\Delta$  in the observed spectrum [13].
  - [16] In Fig.1 of Jido et. al. the state labeled as 1770 and 1660 ought to be turned upside down.
  - [17] The slight difference from Ref. [9] is caused by our definition that all the basis  $\Delta_{1,2}$  and  $N_{1,2}$  have positive parity.